# Orthogonal U(1)'s, Proton Stability and Extra Dimensions

Alfredo Aranda[*] and Christopher D. Carone[†]
*Nuclear and Particle Theory Group*
*Department of Physics, College of William and Mary*
*Williamsburg, VA 23187*

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## Abstract

In models with a low quantum gravity scale, one might expect that all operators consistent with gauge symmetries are present in the low-energy effective theory. If this is the case, some mechanism must be present to adequately suppress operators that violate baryon number. Here we explore the possibility that the desired suppression is a consequence of an additional, spontaneously-broken, non-anomalous U(1) symmetry that is orthogonal to hypercharge. We show that successful models can be constructed in which the additional particle content necessary to cancel anomalies is minimal, and compatible with the constraints from precision electroweak measurements and gauge unification. If unification is sacrificed, and only the new U(1) and its associated Higgs fields live in the bulk, it is possible that the gauge field zero mode and first few Kaluza-Klein excitations lie within the kinematic reach of the Tevatron. For gauge couplings not much smaller than that of hypercharge, we show that these highly leptophobic states could evade detection at Run I, but be discovered at Run II. Our scenario presents an alternative to the 'cartographic' solution to baryon number violation in which leptons and quarks are separated in an extra dimension.

[*]fefo@physics.wm.edu
[†]carone@physics.wm.edu

# 1   Introduction

It is a general principle of effective field theory that one should include all operators consistent with symmetry constraints when constructing a low-energy effective Lagrangian [1]. Such operators are suppressed by powers of the ultraviolet cutoff, so that each has the appropriate mass dimension, and multiplied by coefficients that parameterize the unknown physics relevant at higher energy scales. When this approach is applied to models with a low quantum gravity scale [2], one obtains a multitude of phenomenological disasters, unless specific mechanisms are invoked to suppress contributions to processes that are suppressed or absent in the standard model [3]. In this paper, we consider the possibility that baryon-number-violating operators are present generically in such theories [4], but are suppressed by an additional, non-anomalous, spontaneously-broken U(1) gauge symmetry that is orthogonal to hypercharge [5]. We will argue that the natural scale for the breaking of this symmetry is $\mathcal{O}(1)$ TeV, so that our scenario may have testable consequences at the Fermilab Tevatron, or at the next generation of collider experiments.

We focus on baryon number violation since it is by far the most dangerous of nonstandard model processes. Even if the Planck scale has its conventional value $M_{Pl} \approx 10^{19}$ GeV, the most general set of Planck-suppressed, baryon-number-violating operators lead to proton decay at a rate that is much too fast, unless there is some additional parametric suppression. For example, the superpotential operator $(Q_1 Q_{1,2}) Q_2 L_i / M_{Pl}$ must be suppressed by an additional factor of $\mathcal{O}(10^{-6})$ to avoid conflict with the proton lifetime bounds from SuperKamiokande [6]. For a high Planck scale, this additional suppression factor can originate from the same sequential breaking of flavor symmetries that may account for the smallness of the Yukawa couplings of the first two generations [7]. However, if $M_{Pl}$ is in the $1 - 100$ TeV range, which can be the case in models with extra spacetime dimensions compactified at the TeV-scale, then a much higher degree of suppression is required. We will show that a flavor-universal U(1) gauge symmetry, isomorphic to baryon number on the standard model particle content and spontaneously broken only slightly above the weak scale, is sufficient to avoid any phenomenological problems stemming from baryon-number-violating operators.

It is worth stressing that there are probably many possible ways of suppressing or eliminating proton decay in theories with a low Planck scale. One elegant suggestion made by Arkani-Hamed and Schmaltz is that quarks and leptons may be localized at different points in an extra dimension, so that proton decay operators are suppressed by the tiny overlap of the quark and lepton wave functions [8]. The approach that we consider here is complementary in that it applies also to the case when quarks and leptons are fixed to a single brane, with no separation. No doubt, this possibility has met considerable interest in the recent literature [9].

There is some relationship between the present work and earlier papers on the possibility of gauged baryon number, in which the scale of spontaneous symmetry breaking was taken *below* $M_Z$ [5, 10, 11, 12]. While the proton decay issue was discussed in Ref. [5], the model used as a basis for the argument is now excluded at above the 95% confidence level from

bounds on the electroweak $S$ parameter – the model required a fourth chiral generation to cancel gauge anomalies. Other possibilities for anomaly cancellation discussed in the first version of Ref. [10] are excluded by $S$, and are also inconsistent with gauge coupling unification. Here we will present a supersymmetric model that is consistent with unification (in the case where all gauge and Higgs fields live in the bulk [13, 14]) as well as the anomaly-cancellation constraints. The required extra matter is chiral under the full gauge group, but vector-like under the standard model gauge factors, so that the $S$ parameter bound may be avoided. The extra matter fields get masses of order the U(1) breaking scale $\Lambda_B$, which in principle could be decoupled from the weak scale. We suggest, however, that a natural possibility for generating $\Lambda_B$ is a radiative breaking scenario that relates this scale to the scale of supersymmetry breaking. In this case, the new physics we introduce becomes relevant for TeV-scale collider experiments.

One of the distinctive features of the $Z'$ boson in the class of models we consider is its natural leptophobia. While it may be tempting to think that a model with gauged baryon number is leptophobic by design, it is not hard to see that this statement is patently false. Generically, any additional U(1) symmetry will mix with hypercharge via the kinetic interaction

$$\mathcal{L} = -\frac{1}{2} c_B F_Y^{\mu\nu} F_{\mu\nu}^{new} \ , \tag{1.1}$$

which is not forbidden by any symmetry of the low-energy theory. Even if $c_B$ is identically zero at the ultraviolet cutoff of the theory $M_{Pl}$, it will be renormalized at one loop by all particles that carry both hypercharge and the additional U(1) charge, so that $c_B(\mu) \neq 0$ for $\mu < M_{Pl}$. The class of models that we consider here have the property that $c_B(M_{Pl}) = 0$, and in addition

$$Tr(BY) = 0 \ , \tag{1.2}$$

where $B$ and $Y$ are the baryon number and hypercharge matrices, and the trace sums over all fields in the theory. It is in this sense that we say the additional U(1) is orthogonal to hypercharge. Such orthogonal U(1)'s are known to arise in string theory [15], though we will not commit ourselves to any specific string-theoretic embedding. The constraint $Tr(BY) = 0$ assures that the mixing parameter $c_B(\mu)$ remains zero until the heaviest particle threshold is crossed. In our models, the heaviest particle threshold includes all the nonstandard particles introduced to cancel anomalies; thus the running of $c_B(\mu)$ begins after the exotic states are integrated out, and hence is controlled *solely* by the standard model particle content. This gives our phenomenological analysis a high degree of model independence: a similar model with different nonstandard matter content would have identical $Z'$ phenomenology.[‡]

It is worth stressing that the leptophobia of the $Z'$ in this model (as well as the leptophobia of its Kaluza-Klein excitations) is quite robust. For example, one might think that the $Z'$ could be made less leptophobic by taking the scale $\Lambda_B$ to be high (so that $c_B(\mu)$ would have a greater distance to run). However, this possibility is inconsistent with the

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[‡]For $Z'$ models that suppress proton decay and have a different phenomenology, see Ref. [16].

assumption that (a) the $Z'$ zero mode is phenomenologically relevant and (b) the model is consistent with unification. Since we don't know the string normalization of the new U(1) gauge coupling, we only require that it not differ wildly in strength from hypercharge at low energies. For a $Z'$ with mass $M_B < 1$ TeV, and coupling $g_B \lesssim g_Y$, the associated symmetry breaking scale $M_B/g_B$ cannot be arbitrarily high. Since this is also the scale of the exotic matter content, $c_B(\mu)$ cannot run over very large intervals. If one takes $g_B$ to be smaller, the scale at which running begins is pushed up, but $c_B(\mu)$ runs more slowly due to the reduced coupling. We study this effect quantitatively in Section 3.

Finally, if one is willing to sacrifice simple power-law unification, as in the original scenario of Arkani-Hamed, Dimopoulos and Dvali [2], then it is possible to consider a scenario where only gravity and the additional U(1) may propagate into the extra dimensional bulk space. What is interesting about this possibility is that strongest bounds on the compactification scale come solely from the effects of the new U(1). As a consequence, the $Z'$ and its Kaluza-Klein (KK) excitations may be brought within the kinematic reach of the Tevatron. We show that for gauge couplings not much smaller than that of hypercharge, the $Z'$ and its first few KK modes could remain invisible at Run I of the Tevatron, but be discerned easily at Run II. For this model, the ability of a collider experiment to probe weak couplings is as important as mass reach; we show that the enhanced luminosity of Run II could allow the Tevatron to probe a significant region of the model's parameter space.

In the next section, we highlight the points discussed above by presenting a concrete example. We do not view this model as unique, but rather as a representative example of a class of orthogonal U(1) models that have similar low-energy physics. In Section 3 we discuss the low-energy phenomenology of our scenario, and in the final section present our conclusions.

## 2   A Model

The gauge group is that of the standard model with an additional U(1) factor:

$$G = SU(3) \times SU(2) \times U(1)_Y \times U(1)_B \ . \tag{2.1}$$

We normalize the gauge coupling $g_B$ such that all standard model quarks have charge 1/3, while all leptons and standard model Higgs fields have charge 0; these are the conventional charge assignments for baryon number in the standard model. Gauging this symmetry requires the introduction of exotic matter to cancel chiral gauge anomalies, as well as additional Higgs fields to spontaneously break the symmetry and avoid long-range forces. The aim of this section is to show that this can be done in a relatively simple way, consistent with a number of important phenomenological constraints. In particular, we show that exotic matter can be chosen such that the model (1) is consistent with gauge unification, (2) is anomaly free, (3) suppresses proton decay sufficiently, (4) has no unwanted stable

colored or charged states, and (5) has a mechanism for giving the exotic matter mass. We present the model by considering these issues systematically:

*Gauge Unification.* We would like our model to be consistent with power-law unification [14], at least in the case where all the gauge and Higgs fields are allowed to propagate into the extra-dimensional space. Since the string normalization of the additional U(1) is uncertain [18], we seek to preserve unification of the ordinary standard model gauge factors while allowing $g_B$ to assume values at low energies that do not differ wildly from that of hypercharge. We therefore require that the exotic matter fields fall in complete SU(5) representations. While there are of course other possibilities [17], this is the simplest. We introduce an extra generation that is vector-like under the standard model gauge factors but chiral under U(1)$_B$:

$$
\left. \begin{array}{c} Q_L \\ U_R \\ D_R \end{array} \right\} b_Q \qquad \left. \begin{array}{c} \overline{Q}_R \\ \overline{U}_L \\ \overline{D}_L \end{array} \right\} b_{\bar{Q}}
$$

$$
\left. \begin{array}{c} L_L \\ E_R \\ N_R \end{array} \right\} b_L \qquad \left. \begin{array}{c} \overline{L}_R \\ \overline{E}_L \\ \overline{N}_L \end{array} \right\} b_{\bar{L}}
$$

$$(2.2)$$

Although we assume supersymmetry, we show only the fermionic components above. The overlines indicate Dirac adjoints, and the $b$'s represent the U(1)$_B$ charges, yet to be specified. (Four distinct U(1)$_B$ charges is the smallest number we found that could produce a viable model.) The charges under the standard model gauge factors for fields in the first column are precisely the same as those of fields in an ordinary standard model generation; the only exception is $N_R$ which is a standard model singlet. The fields in the second column have conjugate standard model charges so that, for example, $\overline{Q}_R Q_L$ would be invariant if $b_Q + b_{\bar{Q}} = 0$. As we will see below, our choices for the $b_i$ are such that all the fields in Eq. (2.2) obtain masses of order the U(1)$_B$ breaking scale.

*Anomaly Cancellation* We now aim to restrict the $b_i$ so that the model is free of gauge anomalies. We first note that triangle diagrams involving only standard model gauge factors remain vanishing since the additional matter is introduced in complete generations. We therefore must consider anomalies of the form U(1)$_B^3$, $G_{SM}$U(1)$_B^2$ and $G_{SM}^2$U(1)$_B$, where $G_{SM}$ represents any of the standard model group factors. Given the tracelessness of the non-Abelian generators, this reduces the relevant anomalies to the set: U(1)$_Y$U(1)$_B^2$, SU(3)$^2$U(1)$_B$, SU(2)$^2$U(1)$_B$, U(1)$_Y^2$U(1)$_B$, and U(1)$_B^3$. It is easy to see that the SU(3)$^2$U(1)$_B$ anomaly vanishes since all colored matter with the same U(1)$_B$ charge comes in groups with equal numbers of left- and right-handed fields. The same can be said of the U(1)$_B^3$ anomaly, since the additional $N_{L,R}$ states assure that the exotic 'leptons' with the same U(1)$_B$ charge again come in equal numbers of left- and right-handed fields. Finally, we can dispense with the U(1)$_Y$U(1)$_B^2$ anomaly by noting that every group of particles with the same U(1)$_B$ charge separately satisfies $Tr(Y) = 0$. The remaining two anomaly cancellation conditions,

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$SU(2)^2 U(1)_B$ and $U(1)_Y^2 U(1)_B$, give exactly the same constraint

$$3\Delta_Q + \Delta_L = -3 \ , \tag{2.3}$$

where we have defined

$$\Delta_Q = b_Q + b_{\bar{Q}} \quad \text{and} \quad \Delta_L = b_L + b_{\bar{L}} \ . \tag{2.4}$$

Given the charges defined in Eq. (2.2), we impose Eq. (2.3) to render our theory free of anomalies.

Notice that $-\Delta_Q$ and $-\Delta_L$ also represent the charges of Higgs fields that we require to give the exotic matter fields masses when $U(1)_B$ is spontaneously broken. The most economical exotic Higgs sector is obtained by setting

$$\Delta_Q = \pm\Delta_L \ . \tag{2.5}$$

Then all the desired mass terms may be formed by introducing a single pair of Higgs fields

$$S_B \quad \text{and} \quad S_{\bar{B}} \ , \tag{2.6}$$

with charges $+\Delta_Q$ and $-\Delta_Q$, respectively. This is the minimal possibility, since, as in the minimal supersymmetric standard model (MSSM), a vector-like pair of Higgs superfields is required to avoid additional anomalies. The choice of Eq. (2.5) together with the constraint Eq. (2.3) implies that either

$$\Delta_Q = \Delta_L = -3/4 \quad \text{or} \quad \Delta_Q = -\Delta_L = -3/2 \ . \tag{2.7}$$

The remaining freedom to choose exotic $U(1)_B$ charges will be important in satisfying the other phenomenological constraints below.

*Proton decay.* If our additional U(1) symmetry were unbroken, then it would be clear that all operators contributing to proton decay would be exactly forbidden. When the symmetry is spontaneously broken, the form of baryon-number-violating operators in the low-energy effective theory depends on the charge assignment of the Higgs fields which break $U(1)_B$, as well as on the size of their vacuum expectation values (vevs). Let us work in the very low-energy limit, below the scales of extra dimensions, exotic matter, and supersymmetry breaking, which we will take to be $\sim 1$ TeV universally for the purposes of the present argument. In this effective nonsupersymmetric theory, operators that could contribute to proton decay have the form [5]

$$\mathcal{O} = q^k \ell^m \chi^n \ , \tag{2.8}$$

where $q$ and $\ell$ represent generic quark and lepton fields, respectively, and $\chi$ represents the vev of either $S_B$ or $S_{\bar{B}}$. Here we have suppressed both the Dirac structure of the operator and the standard model gauge indices for convenience. First, we note that since the lepton electric charge is integral, $k$ must be a multiple of 3, *i.e.* $k = 3p$. It follows that the baryon

number of $q^k \equiv q^{3p}$ is $p$, which is an integer. On the other hand, this must be compensated by the baryon number of $\chi$, which is either $\pm 3/2$ or $\pm 3/4$, given the charges of the $S_B$ fields already discussed. Thus we conclude that the operators represented by Eq. (2.8) must be of the form

$$(q^9 \chi^2)^r \ell^m \quad \text{or} \quad (q^9 \chi^4)^r \ell^m \;, \tag{2.9}$$

where $r$ and $m$ are integers. The point is simple: the fact that the possible symmetry breaking 'spurions' have fractional $U(1)_B$ charges forces the baryon-number-violating operators to contribute to no less than $\Delta B = 3$ transitions. This renders our model safe from proton decay as well as $N$-$\overline{N}$ oscillations. The operators in Eq. (2.9) are suppressed by high powers of mass scales that are either 1 TeV or $M_{Pl}$, and thus are unlikely to have any observable effects on stable matter at low energies.

*Avoiding Stable Charged Exotic Matter.* We will now further restrict our charge assignments $b_i$ to assure that we have no stable heavy states that are charged under any of the standard model gauge factors. This allows us to evade bounds on stable charged matter from searches for anomalously heavy isotopes in sea water [19]. In both the exotic lepton and quark sectors separately, it is always possible to choose Yukawa couplings such that one exotic state is lightest, and ordinary weak decays to this state are kinematically allowed. For example, the exotic lepton superpotential couplings (in terms of left-handed chiral superfields)

$$W \supset L\bar{L}S_{\bar{B}} + (E\bar{E} + N\overline{N})S_B + (LE + \bar{L}\bar{N})H_D + (\bar{L}\bar{E} + LN)H_U \tag{2.10}$$

lead to mass terms of the form

$$( \; \bar{e}_H \quad E \; ) \begin{pmatrix} M_1 & m_2 \\ m_1 & M_2 \end{pmatrix} \begin{pmatrix} e_H \\ \bar{E} \end{pmatrix} + ( \; \bar{\nu}_H \quad N \; ) \begin{pmatrix} M_1 & m_4 \\ m_3 & M_3 \end{pmatrix} \begin{pmatrix} \nu_H \\ \overline{N} \end{pmatrix} \;, \tag{2.11}$$

where the $M_i$ are masses of order the $U(1)_B$ breaking scale, while $m_i$ are of order the weak scale. Here we have written the component superfields in the doublets $L$ ($\overline{L}$) as $\nu_H$ ($\overline{\nu}_H$) and $e_H$ ($\overline{e}_H$). Clearly one has the freedom to arrange for the lightest exotic lepton state to be neutral. For example, for the specific choice $M_1 = M_2 = M_3$, $m_1 = m_2$ and $m_3 = m_4$, the lightest charged state has mass $M - m_1$ while the lightest neutral state $M - m_3$; we therefore could take $m_1 < m_3$. In the exotic quark sector, the lightest state is charged and colored, so some additional mechanism must be provided to assure it decays to ordinary particles. Since we are working in the context of models in which the Planck scale is low, we can make the lightest exotic quark unstable by considering possible higher-dimension operators, allowed by the symmetries of the theory and suppressed by the cutoff. As there is some freedom in how we may accomplish this, let us restrict our subsequent discussion to a specific example. Let us choose the charge assignment in which $\Delta_Q = -3/2$. The choice $b_Q = -2/3$ and $b_{\bar{Q}} = -5/6$ is consistent with this condition, and also allows the superpotential operator

$$\frac{1}{M_{Pl}} q\, q\, Q\, \ell \;, \tag{2.12}$$

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where lower-case superfields are those of the standard model. This operator allows for three-body decays for the lightest exotic quark field (for example, to a normal lepton and two squarks). Even if the superpartners are heavy, so that this decay is not kinematically allowed, one can obtain a four-fermion operator by "dressing" Eq. (2.12) with a gaugino exchange. In this case, the decay proceeds to two quarks and a lepton, with a width of order

$$\Gamma \sim \frac{1}{64\pi^3} \left(\frac{1}{16\pi^2}\right)^2 \left(\frac{M_Q}{M_{Pl}}\right)^2 M_Q \ . \tag{2.13}$$

The first factor is from three-body phase space, the second from the fact that the amplitude occurs at one-loop, and the rest follows from dimensional analysis. The lightest exotic quark decays well before nucleosynthesis providing that $M_{Pl} < 10^{13}$ GeV; this is not a problem in our scenario. Note that the charge assignments $b_Q = -2/3$ and $b_{\bar{Q}} = -5/6$ assure that potentially dangerous mass mixing terms like $q\overline{Q}$, and $QH_D d$ have $U(1)_B$ charges of $-1/2$ and $-1$, respectively. Since this is not an integral multiple of $3/2$ (the magnitude of the exotic Higgs' $U(1)_B$ charges) such operators are forbidden by the gauge symmetry. We will adopt the present choice of $b_Q$ and $b_{\bar{Q}}$ for the subsequent discussion. However, the reader should keep in mind that other possible assignments may render the exotic matter unstable, given the presence of higher-dimension operators at the relatively low cutoff of the theory.

*Orthogonality.* The only charges we have not yet fixed are $b_L$ and $b_{\bar{L}}$, which have been constrained such that $b_L + b_{\bar{L}} = 3/2$. Since we wish to restrict our discussion to models that satisfy $Tr(BY) = 0$, we fix our remaining degree of freedom by imposing this constraint. It is straightforward to check that $Tr(BY) = 9 \cdot \frac{1}{3} \cdot (2 \cdot \frac{1}{6} + \frac{2}{3} - \frac{1}{3}) = 2$ for the ordinary matter, where the overall factor of 9 is the multiplicity due to color and number of generations. For the exotic matter, the quark fields contribute $Tr(BY) = 3 \cdot (b_Q - b_{\bar{Q}}) \cdot (2 \cdot \frac{1}{6} + \frac{2}{3} - \frac{1}{3}) = 1/3$ given our previous choice of $b_Q = -2/3$ and $b_{\bar{Q}} = -5/6$. We now choose $b_L = 4/3$ and $b_{\bar{L}} = 1/6$. The exotic lepton contribution is then $Tr(BY) = (b_L - b_{\bar{L}})(2 \cdot [-\frac{1}{2}] - 1) = -7/3$. Hence, the orthogonality of $U(1)_B$ and hypercharge is maintained. Notice that our choice for $b_L$ and $b_{\bar{L}}$ is such that no dangerous mass mixing terms between exotic and standard model leptons are generated after $U(1)_B$ is spontaneously broken. Now that all our charges have been fixed, we summarize them here for convenience:

$$\begin{array}{ll} b_Q = -2/3 & b_{\bar{Q}} = -5/6 \\ b_L = \phantom{-}4/3 & b_{\bar{L}} = \phantom{-}1/6 \end{array} \ . \tag{2.14}$$

*Symmetry Breaking.* It is customary in model building to avoid discussing the origin of symmetry breaking scales, given the model-dependence that this issue often entails. Here we only aim to emphasize that the scale of $U(1)_B$ breaking may be tied quite naturally to the scale of supersymmetry breaking. This point is worth mentioning given that we have constructed our model specifically to allow for the decoupling of the nonstandard sector, to avoid bounds from precision electroweak measurements. One way in which the supersymmetry breaking and $U(1)_B$ scale may be related is if the potential for the nonstandard

Higgs fields $S_B$ and $S_{\bar{B}}$ develops its vacuum expectation value as a consequence of a soft scalar squared mass running negative, the analog of the radiative breaking scenario in the MSSM. This scenario can be implemented in the present context since the exotic Higgs fields couple to a sector of new matter fields with large Yukawa couplings. The exotic Higgs fields have the superpotential coupling

$$W = \mu_s S_B S_{\bar{B}}, \tag{2.15}$$

the analog of the $\mu$ term in the MSSM. Introducing soft supersymmetry breaking masses, and $D$-terms, the scalar potential for the exotic Higgs fields is given by

$$
\begin{aligned}
V &= \frac{1}{2}(\mu_s^2 + m_B^2)(s_B^2 + p_B^2) + \frac{1}{2}(\mu_s^2 + m_{\bar{B}}^2)(s_{\bar{B}}^2 + p_{\bar{B}}^2) \\
&+ \mu_s B_s(s_B s_{\bar{B}} - p_B p_{\bar{B}}) + \frac{9}{32}g_B^2(s_B^2 + p_B^2 - s_{\bar{B}}^2 - p_{\bar{B}}^2)^2 \ ,
\end{aligned} \tag{2.16}
$$

where $s_{B,\bar{B}}$ and $p_{B,\bar{B}}$ represent the scalar and pseudoscalar components of each of the fields, and $m_B$, $m_{\bar{B}}$, and $B_s$ are soft, supersymmetry-breaking masses. It is straightforward to show that this potential has stable (local) minima in which one scalar squared mass is negative and both $S_B$ and $S_{\bar{B}}$ acquire vacuum expectation values. For example, for the parameter choice $g_B = 0.3$, $\mu_s = 1$ TeV $B_s = -1$ TeV, $m_B^2 = -1.48$ TeV$^2$, and $m_{\bar{B}}^2 = 2.81$ TeV$^2$, we find the vevs

$$\langle s_B \rangle = 3 \text{ TeV} \quad \langle s_{\bar{B}} \rangle = 1 \text{ TeV},$$

the scalar squared masses
$$0.99 \text{ TeV}^2 \quad 4.37 \text{ TeV}^2 \ ,$$

and the pseudoscalar squared mass

$$3.33 \text{ TeV}^2 \ .$$

These are acceptable values. Another possible form for the potential is that of the next-to-minimal supersymmetric standard model, in which both the ordinary $\mu$ parameter and the parameter $\mu_s$ could have a common origin, the vev of a singlet field. We will not study the issue of possible potentials any further here, though such an investigation would be required if experimental evidence for the model became available.

# 3 Phenomenology

In this section, we explore the $Z'$ phenomenology of our model. We will assume for simplicity that the scale of exotic matter, $\Lambda_B$, and of superpartner masses is 1 TeV. The compactification scale, which we call $\Lambda$ below, is a free parameter. In the case where all non-chiral matter (i.e. the Higgs and gauge fields) are allowed to propagate in the bulk,

we require $\Lambda$ to be greater than a few TeV, to satisfy the constraints from precision electroweak measurements [20]. In this case, the phenomenology that we study is that of the new zero mode gauge field. However, we will also consider the (non-unifiable) possibility that only U(1)$_B$ lives in the bulk, in which case the bounds on $\Lambda$ are substantially weakened. For this choice, the $Z'$ zero mode and first few KK excitations become relevant at planned collider experiments, and will be the focus of our discussion. For concreteness, we perform our numerical analysis in the case of one extra dimension.[§] For more than one extra dimension, the sums involving the KK modes are divergent and must be regulated by some additional, string-theoretic mechanism. We restrict ourselves to one extra dimension to avoid this model-dependent issue; however, the reader should keep in mind that our bounds on the U(1)$_B$ KK modes may be overestimates if there is a mechanism, e.g. brane recoil effects [21], that suppresses the KK couplings.

One of the interesting properties of this class of models, regardless of which case we consider, is the strong leptophobia of the $Z'$ and its KK excitations. Given our assumption of a vanishing kinetic mixing parameter, $c_B$, at the string scale, $c_B$ remains vanishing down to the scale of exotic matter, since $Tr(BY) = 0$. At lower scales, the exotic states are integrated out of the theory, and the orthogonality constraint is no longer satisfied. With our choice of energy scales, $c_B$ remains small down to the $Z'$ mass, so we may treat Eq. (1.1) as a perturbative interaction. Thus, the Feynman rule for the $Z'$-hypercharge vertex is given by

$$- i\, c_B \left( p^2 g^{\mu\nu} - p^\mu p^\nu \right) \, . \tag{3.1}$$

Since we assume that the scale of superpartner masses is the same as the scale of exotic matter, we evaluate the non-supersymmetric running of $c_B$; at one-loop we obtain the renormalization group equation (RGE)

$$\mu \frac{\partial}{\partial \mu} c_B = -\frac{1}{3\pi} \sqrt{\alpha_Y \alpha_B} \left[ \frac{5}{6} N_u - \frac{1}{6} N_d \right] \, , \tag{3.2}$$

where $N_u$ and $N_d$ are the number of standard model up-type and down-type quarks propagating in the loop. This RGE is solved subject to the boundary condition $c_B(\Lambda_B) = 0$, for the reasons described above. Notice that the running of $c_B$ is controlled entirely by the standard model particle content, since these are the only fields relevant below the scale $\Lambda_B$. Thus, our analysis is independent of the specific exotic sector introduced to cancel anomalies.

We may now consider the phenomenology of the model by determining bounds in the $M_B$-$\alpha_B$ plane. We will assume $M_B > m_{top}$ (which was not studied in Refs. [5, 10]) and first consider the case in which all non-chiral superfields live in the bulk. For most of the mass range of interest, the $Z'$ will be sufficiently heavier than the $Z$ so that the most stringent bounds are obtained from direct collider searches. We consider the limits on $Z'$'s decaying to dijets and dileptons at the Fermilab Tevatron Collider:

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[§]Of course, gravity also lives in the bulk. Our model does not preclude the possibility that gravity propagates in a larger number of dimensions than U(1)$_B$.
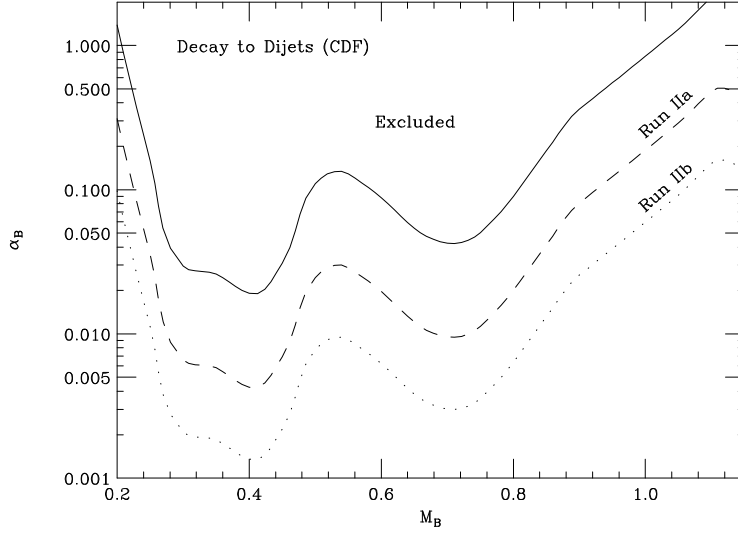
Figure 1: Bound on $\alpha_B$ from the cross section times branching fraction to dijets. The solid line corresponds to the bound obtained from Run I with a Luminosity of 106 pb$^{-1}$. The dashed line corresponds to a luminosity of 2 fb$^{-1}$ for Run IIa and the dotted line to a luminosity of 20 fb$^{-1}$ for Run IIb.

*Decays to Dijets.* The CDF Collaboration has placed bounds on narrow resonances decaying to dijets in $p\bar{p}$ collisions at $\sqrt{s} = 1.8$ TeV [22]. They present the 95% C.L. upper limits on cross section times branching ratio as a function of the $Z'$ mass in the range $0.2 - 1.15$ TeV. Since the kinetic mixing effects are small (as we will see below), the branching fraction to dijets in our model is nearly 100%; thus we compare the CDF bounds to the $Z'$ production cross section in our model, which we estimate using the narrow width approximation:

$$\sigma(p\bar{p} \to Z' \to dijets) = \frac{4\pi^2}{9} \frac{\alpha_B}{s} \int dy \sum_{i,j} f_i^p(y, \sqrt{s}, M_B) f_j^{\bar{p}}(y, \sqrt{s}, M_B) \,. \qquad (3.3)$$

Here $y$ is the rapidity, $\sqrt{s}$ is the center of mass energy, and $f^p$ ($f^{\bar{p}}$) represents the appropriate parton distribution functions for $p\bar{p}$ collisions. Using the CTEQ 4M structure functions [23] at $\sqrt{s} = 1.8$ TeV for our numerical analysis, we obtain a bound on $\alpha_B(M_B)$ as a function of $M_B$, shown in Fig. 1. The solid line corresponds to the Run I luminosity ($\mathcal{L}$) of $\sim 0.1$ fb$^{-1}$ and is the strongest bound on the model. We also estimate the ability of the Tevatron to probe additional parameter space at Run II. Note that the shape of the excluded region in Fig. 1 depends on a detailed analysis of both statistical and experimental systematic uncertainties; the latter are difficult to extrapolate with precision to Run II. Therefore, we rely instead on the observation that statistical and systematic uncertainties generally both scale as $\sqrt{\mathcal{L}}$ (*i.e.* the systematic uncertainties can be reduced by higher statistics). Thus, we make a simple extrapolation, scaling the bound from Run I down by $\sqrt{\mathcal{L}_I}/\sqrt{\mathcal{L}_{II}}$ using
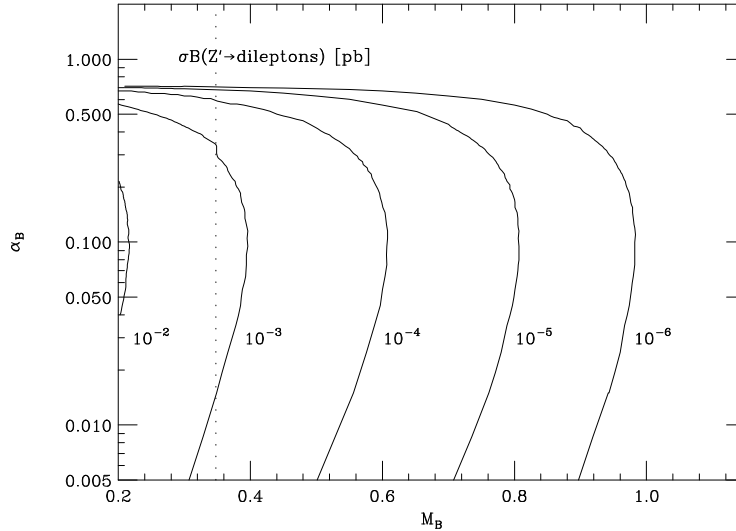
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Figure 2: Contours of constant cross section times branching fraction to dileptons. The dotted line shows the threshold $M_B = 2m_{top}$.

the expected luminosities at Run IIa and Run IIb, 2 and 20 fb$^{-1}$ respectively; this yields the two other curves shown in Fig. 1. We see that, for example, it is possible to have a new gauge boson in the region between 500 and 600 GeV with a coupling of electromagnetic strength that could be observed at Run II.

*Decays to Dileptons.* Given the construction of our model, the specification of $M_B$ and $\alpha_B$ is sufficient to determine the magnitude of $c_B(M_B)$, up to a small uncertainty. For each point in the parameter space, $M_B/\sqrt{4\pi\alpha_B}$ is of order the scale of U(1)$_B$ breaking. However, this scale also determines the masses of the exotic fermions, and the point at which $c_B$ begins to run. The only uncertainty is in the Yukawa couplings of the exotic matter, which we assume is of order one (say, between 1/3 and 3); this only affects the result logarithmically. To account for the mixing, we use Eq. (3.2) to run $c_B$ from the U(1)$_B$ breaking scale $\Lambda_B = rM_B/g_B$, where $r$ is an $\mathcal{O}(1)$ uncertainty, down to $M_B$ with the condition $c_B(\Lambda_B) = 0$. We show some typical values of $c_B$ in Table 1 for different choices of $M_B$ and $\alpha_B$. The results are uniformly small, due to two competing effects: if the coupling $g_B$ is reduced with $M_B$ held fixed, then the 'starting' scale $\Lambda_B$ is increased, while the rate of running, *i.e.* the right-hand side of Eq. (3.2), is reduced. As a consequence, the branching fraction to leptons

$$B = \frac{\frac{3}{2}c_B^2\alpha_Y}{\frac{N_f}{9}\alpha_B + \frac{3}{2}c_B^2\alpha_Y} \, , \tag{3.4}$$

is highly suppressed throughout the parameter space in Fig. 1. Here $N_f$ is the number of quarks lighter than $M_B/2$. In Fig. 2 we show contours of constant $\sigma$B; note that $\sigma$B vanishes when $\Lambda_B(\alpha_B) = M_B$. The CDF bound on this product in no stronger than 0.04 pb for dilepton invariant masses above $\sim 400$ GeV, and is significantly weaker for
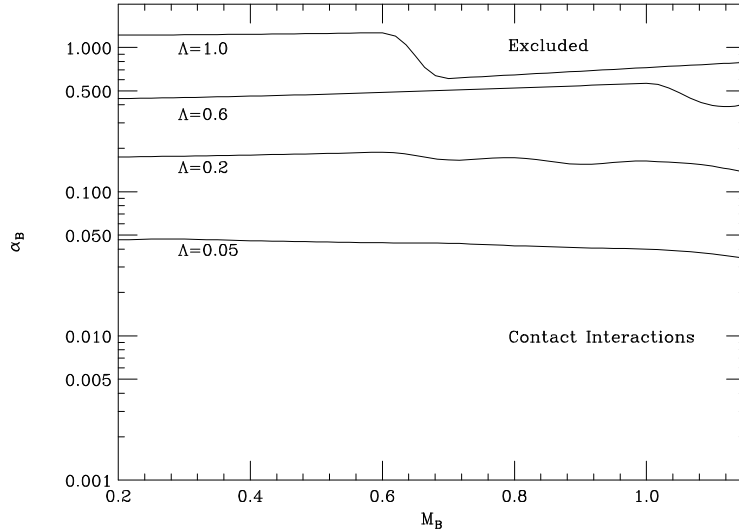
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Figure 3: Bounds obtained from the contribution of KK modes heavier than 1.15 TeV to contact interactions for several values of $\Lambda$.

smaller masses [24]; as a consequence, no additional bound can be placed on our parameter space. It is possible, however, that a dilepton signal could be discerned at Run II, if the $Z'$ were already discovered in the dijet channel. For example, for $M_B \approx 400$ GeV, where the current bound is 0.04 pb, a simple rescaling by $\mathcal{L}$ suggests that the bound could become 0.0028 pb after 20 fb$^{-1}$ of integrated luminosity. The results in Fig. 2 imply that this would be sufficient to see the model's tiny dilepton signal.

*KK-modes.* The $Z'$ phenomenology we have discussed thus far has related to the zero-mode gauge field, and is independent of how the model is configured in extra dimensions. As we mentioned earlier, if all the non-chiral fields propagate in the bulk, then the first $Z'$ KK mode is outside the reach of the Tevatron, and the zero-mode is of principle interest to us. Here, we wish to consider an alternative possibility, that the compactification scale is low enough such that the first few KK modes are also within the kinematic reach of the

| $M_B$ (TeV) | $\alpha_B(M_B)$ | $c_B(M_B)$ |
|:---:|:---:|:---:|
| 0.2 | 0.1 | 0.00688 |
| 0.5 | 0.1 | 0.00694 |
| 1.0 | 0.1 | 0.00699 |
| 0.2 | 0.01 | 0.00469 |
| 0.5 | 0.01 | 0.00471 |
| 1.0 | 0.01 | 0.00473 |

Table 1: Kinetic mixing for $r = 3$.

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Tevatron. This can be the case if only $U(1)_B$ and its associated exotic Higgs fields live in the bulk. The usual strong bounds on $\Lambda$ are evaded in this situation since there are *no* exotic Higgs fields charged under both $U(1)_B$ and any of the standard model electroweak gauge factors – the vev of such a field would lead to mixing at tree-level between the $Z$ and $Z'$ KK modes. In order to determine the relevant bounds, let us consider the following terms in the $Z'$ Lagrangian:

$$
\mathcal{L}_{KK} = -\frac{1}{4}\sum_{n=0} F^{(n)}_{\mu\nu} F^{\mu\nu}_{(n)} + \frac{1}{2}\sum_{n=0}\left( M_B^2 + \Lambda^2 n^2\right) Z'^{\mu(n)} Z'^{(n)}_\mu
$$
$$
- \frac{g_B}{3}\,\bar{q}\gamma^\mu\left( Z'^{(0)}_\mu + \sqrt{2}\sum_{n=1} Z'^{(n)}_\mu\right) q\,. \tag{3.5}
$$

Notice that the KK modes have contributions to their masses from both the symmetry breaking and the compactifaction scale. If $\Lambda \ll M_B$, there is effectively a 'pile-up' of states with masses of order $M_B$ and multiplicity $M_B/\Lambda$. This is one way in which low-energy bounds are enhanced. In addition, the coupling of the KK modes to quarks has an extra factor of $\sqrt{2}$ compared to the coupling of the zero mode; this results from the field rescalings necessary to put the four-dimensional kinetic terms in canonical form, and to give the zero-mode gauge coupling its conventional normalization. Hence, the appropriate dijet bound on a given KK mode may be obtained from Fig. 1 by scaling down the exclusion limit shown by a factor of 2.[¶] If $\Lambda$ is sufficiently small, the zero mode and first few KK modes could be unobserved in Run I, but discovered at Run II. We therefore consider whether $\Lambda$ can be small enough for this interesting situation to be obtained.

Aside from the KK modes that are within the reach of the Tevatron, there is also an infinite tower of heavier modes that are integrated out of the low-energy theory. Thus, the new physics manifests itself as a series of narrow resonances, together with effective contact interactions that lead to smoothly growing cross sections. We may use the bounds on four-quark contact interactions to bound the compactification scale. If we integrate out all the modes with mass $M_B > M_{min} = 1.15$ TeV (the endpoint of the dijet invariant mass spectrum in Ref. [22]) we obtain operators of the form

$$
\mathcal{L}_{\bar{q}q\bar{q}q} = -\sum_{n_{min}}^\infty \frac{g_B^2}{9M_n^2}\,\bar{q}_L\gamma_\mu q_L \bar{q}_L\gamma^\mu q_L + \cdots \tag{3.6}
$$

where $M_n^2 = M_B^2 + n^2\Lambda^2$, and $n_{min}$ corresponds to the first KK mode above $M_{min}$. We show only the purely left-handed operator, which is the one most tightly constrained of those listed in the Review of Particle Physics [19], *viz.*, $\Lambda^-_{LL}(qqqq) > 2.4$ TeV at 95% C.L., with $\Lambda^-_{LL}(qqqq)$ defined therein. The sum shown in Eq. (3.6) can be evaluated analytically so that the bound may be written as

$$
\alpha_B < \frac{9M_B\Lambda}{(2.4\text{TeV})^2}\left[ i\Psi(n_{min} - \frac{iM_B}{\Lambda}) - i\Psi(n_{min} + \frac{iM_B}{\Lambda})\right]^{-1} \tag{3.7}
$$

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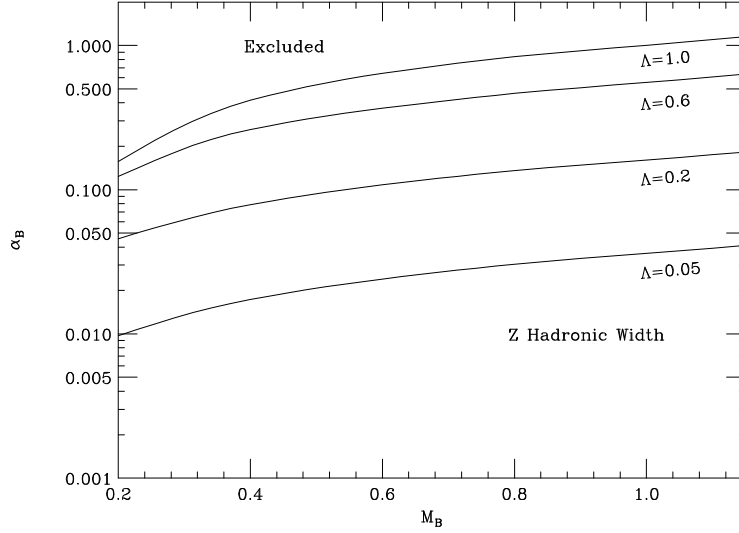[¶]The running of $\alpha_B$ in the range shown in Figure 1 is small, and can be neglected in this discussion.

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Figure 4: Bound obtained from the contribution of the first 1000 KK modes to the Z hadronic width.

where $\Psi(x) = \frac{\partial}{\partial x}[\ln \Gamma(x)]$ is the digamma function. We plot Eq. (3.7) for several values of $\Lambda$ in Fig. 3. The mild steps in these contours occur each time a KK mode becomes more massive than $M_{min}$, and is included in the contact term.

In the case where $\Lambda$ is small, we can also determine whether the pile-up of states at $M_B$ is significantly bounded by $Z$-pole observables. The most stringent constraint for this type of model comes from the measurement of the $Z$ hadronic width [5], which is known to approximately 0.1% [19]. We include contributions from the $Z-Z'$ mixing [5] and from the one-loop $q\bar{q}Z$ vertex correction [10]. The total effect is given by

$$\begin{aligned}\frac{\Delta\Gamma_{had}}{\Gamma_{had}} &\approx -1.194\, c_B(m_Z)\, \sqrt{\alpha_B}\, m_Z^2 \left(\frac{1}{m_Z^2 - M_B^2} + 2\sum_{n=1}^{\infty} \frac{1}{m_Z^2 - M_n^2}\right) \\ &+ \frac{\alpha_B}{18\pi}\left(F_2(M_B) + 2\sum_{n=1}^{\infty} F_2(M_n)\right),\end{aligned} \tag{3.8}$$

where $c_B(m_Z)$ is found by solving Eq. (3.2), and $F_2(M)$ is a loop integral factor that can be found in Ref. [10]. The sums appear linearly in Eq. (3.8) since the effects of new physics appear in an interference term at lowest order. Figure 4 shows the $2\sigma$ bound for several choices of $\Lambda$, where the sum includes the first 1000 KK modes. Generally, the bound obtained from the Z hadronic width supersedes the one obtained from contact interactions. Figs. 3 and 4 in conjunction with Fig. 1 show that the compactification scale $\Lambda$ can be made small enough so that the $Z'$ zero mode and first few KK excitations could be undetectable at Run I and discovered at Run II, without requiring the coupling $\alpha_B$ to be unexplicably small. For example, the parameter choice $\alpha_B = 0.01$, $M_B = 400$ GeV, and $\Lambda = 200$ GeV is consistent with all our constraints.

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# 4 Conclusions

We have shown in this article that it is possible to construct viable models with a non-anomalous U(1) symmetry that is orthogonal to hypercharge and that preserves proton stability, a concern when the quantum gravity scale is low. While exotic chiral fields are required to cancel anomalies, we show that these fields may nonetheless be vector-like under the standard model subgroup, so that constraints from the $S$ parameter are evaded, and may appear in complete SU(5) representations, so that power-law unification may be preserved. The new gauge boson and its KK excitations exhibit a high degree of leptophobia, which is only violated by kinetic mixing with hypercharge, which is small and calculable, given our assumed boundary conditions. If power-law unification is sacrificed, then one may consider the case in which only the extra U(1) lives in the bulk. In this case, the most important bounds on the compactification scale come from processes associated with the exchange of the $Z'$ and its KK excitations, and were found to be relatively weak. This allows the $Z'$ and its first few KK modes to be within the kinematic reach of the Tevatron. In both versions of the model, we considered bounds from collider searches for new particles decaying to dijets and dileptons, and, in the second case, bounds on the compactification scale from contact interactions and contributions to the Z hadronic width. For gauge couplings comparable to that of hypercharge, we showed that this scenario is allowed by current experiments, and that the new gauge boson, and perhaps some of its KK excitations could be discovered by the Tevatron at Run II.

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